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THE CHANGES IN PROPERTIES OF HADRONS IN NUCLEI AND NUCLEAR MATTER*

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ABSTRACT

The changes of hadron properties in dense and/or hot matter are discussed in terms of effective chiral Lagrangians with the parameters of the theory scaled in a simple way. The phenomenologically successful Walecka model is identified as a mean field chiral Lagrangian with the scaled parameters. Kaon condensation and chiral restoration transitions can be described within the same mean field framework.

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Quantum chromodynamics (QCD) tells us that most, if not all, of light hadron masses are generated spontaneously by the breaking of chiral symmetry from $SU(N_f) \times SU(N_f)$ to diagonal $SU(N_f)$ where N_f is the number of flavors, equal to 2 without strangeness and 3 with. It is also widely believed that as a hadronic system is heated to high temperature or compressed to high density, the broken symmetry will get restored in a way paralleling what happens in condensed matter physics. A natural consequence of the restoration of the chiral symmetry must then be that the spontaneously generated masses disappear as density (and perhaps also temperature) is increased. The question we are raising is how does this “shedding of mass” occur?

This question is at the core of the fundamental theory of matter: How is the mass generated, starting with the lightest object like neutrinos to the heaviest detected particle like the top quark?

The aim of this talk is to describe how hadron properties get modified in medium as the system is heated or compressed. That is, immerse a hadron in medium and compress the system or heat it. What does one expect to see happening?

To answer this question, let me start with the simplest nuclear system, namely the deuterium. Let us look at what happens when a soft photon is sent in to probe the system. Consider therefore the well-known inverse process

$$n + p \rightarrow d + \gamma \tag{1}$$

at thermal neutron energy. This process was first explained in a quantitative way by Riska and Brown [1] in 1973. What I will do here is to describe it more accurately in a modern QCD framework. Since it is a very low-energy process, QCD can be represented by an effective chiral Lagrangian field theory. This is because at long wavelength limit, chiral perturbation theory (ChPT) is believed to be exactly equivalent to QCD [2]. This invites us to attempt to describe (1) in terms of a chiral Lagrangian. There is one basic problem in doing this and that has to do with the description of the deuteron in QCD: we do not really know how to derive the deuteron starting *directly* from QCD. Since we are focusing on the chiral aspect of the problem, however, the solution might be sought in a chiral Lagrangian approach to the structure of the deuteron. If the number of colors N_c is in some sense big – which gives rise to what is known as the “large N_c limit” – then the effective Lagrangian is given by meson fields only as we know from the skyrmion structure. There is some important progress in obtaining a bound deuteron, recently through the work of Manton [3], as a baryon number 2 skyrmion, but we are still far from understanding it quantitatively. However there is an indirect approach to this which is consistent with QCD and which has the potential to be quantitatively accurate, namely that as we have known all along, the deuteron is made up of a proton and a neutron bound by meson exchanges:

In the framework of QCD, the nucleon may be gotten from a large N_c Lagrangian as a soliton (skyrmion) but this is now known to be equivalent, at least in the large N_c limit, to having the nucleon as a matter field in the chiral Lagrangian. We are thus led to consider a chiral Lagrangian that contains baryons (nucleon, Δ etc.), pseudo-Goldstone bosons π^i (pions, kaons etc.), vectors $V_\mu = \omega_\mu, \rho_\mu, \dots$ etc. with suitable chiral invariant couplings. There have been some attempts to compute the deuteron from such a Lagrangian in chiral perturbation theory [4] but at present the calculation can be done only at low orders since higher order calculations would involve too many parameters to be completely determined from available experiments. Luckily for our purpose, we need not compute the deuteron from first principles as I shall argue below.

The important point to note is that certain aspects of the deuteron which have to do with the chiral symmetry structure of hadrons can be probed by the process (1) *without knowing* how to get the nucleus itself from a chiral perturbation theory, as recently discussed by Park, Min and Rho [5]. Briefly the argument goes as follows: For physics with energy scale much less than the chiral scale $\Lambda_\chi \sim 4\pi f_\pi \sim 1$ GeV (where f_π is the pion decay constant ~ 93 MeV), the relevant Lagrangian is, schematically,

$$\mathcal{L} = \sum_i \mathcal{L}_i[B, U, \mathcal{M}] \quad (2)$$

where B_i are the baryon fields (both octet and decuplet), U the unitary pseudo-Goldstone fields $U = e^{i\pi/f_\pi}$ and \mathcal{M} the quark mass matrix. The degrees of freedom more massive than the chiral scale Λ_χ are integrated out, appearing implicitly in the counter terms of the Lagrangian. Chiral symmetry requires that there be only derivative couplings apart from terms involving the mass matrix and hence effectively the Lagrangian is an expansion in ∂/Λ_χ and $\mathcal{M}/\Lambda_\chi^2$. Since the baryon mass is $\sim \Lambda_\chi$, the baryons should be introduced as static matter fields, so the derivative on the baryon field does not involve time derivatives. ChPT is then a systematic perturbative expansion in powers of Q , say, Q^n , where Q is the energy-momentum scale being probed, with suitable counter terms to remove divergences and to take account of the degrees of freedom that are integrated out. Now in applying this theory to nuclear systems, we need to separate the class of Feynman diagrams into two, one “irreducible” and the other “reducible.” It is in calculating the irreducible diagrams that ChPT enters. The reducible diagrams – that cannot be treated by ChPT because of infrared singularity – are incorporated by solving a Lippman-Schwinger equation or Schrödinger equation with the potential obtained with the irreducible graphs by ChPT. This is how bound states are to be treated in ChPT. Now in calculating the process (1), we can write the EM current in two terms, one the single-particle current $j^{(1)}$ and the other the two-body current $j^{(2)}$

$$j_\mu^{EM} = j_\mu^{(1)} + j_\mu^{(2)}. \quad (3)$$

The former is called “impulse approximation current” and the latter “exchange current.” For the system considered, we terminate with the two-body current. Later we will see that in heavy nuclei there can enter many-body currents, some of which become quite important.

Very accurate wave functions for the final deuteron and the initial neutron-proton system obtained from some accurate phenomenological potential such as the Argonne v_{18} potential [6] (or a potential calculated in a high order ChPT if it is feasible) would correspond to a high order chiral expansion since the Schrödinger equation sums a certain class of chiral series to all orders and presumably the phenomenological potential also subsumes all orders of chiral perturbation. Now the idea is to compute the matrix element of the current (3) in chiral perturbation expansion in such a way that is consistent with the calculation of the wave functions. The calculation of the one-body current $j^{(1)}$ is without ambiguity. That of the two-body exchange current is somewhat subtle, requiring a careful sorting of irreducible and reducible contributions such that the reducible ones are suitably accounted for in the one-body term with the accurate wave functions. In this way, the Ward identities associated with the conserved vector current are satisfied to a given chiral order in the EM current. This is a numerically accurate procedure.

A recent calculation [5] of the process (1) is given in Fig.1. The calculation was done to order Q^3 which corresponds to next-to-next-to leading order in chiral expansion, that is, to one-loop order. Given the nucleon mass m_N , the pion decay constant f_π , the pion mass m_π , the axial coupling constant g_A and the vector meson mass m_V all *determined in free space*, all the parameters that appear in the theory are fixed in the theory except for the hard core radius r_c in the wave function reflecting on our inability to handle in ChPT very short-range physics. The remarkable agreement with the experiment shows that the two-nucleon systems we are looking at are made up of two nucleons with their properties as given in free space: the chiral Lagrangian with the vacuum values of hadron parameters describes nature remarkably well at low chiral orders.

A very valuable information for heavier and denser nuclei is lodged in some of the terms that are *negligible* in the process (1). As I showed elsewhere [7], two-body currents involving four baryon fields (e.g. Fig.2(c) with the blob replaced by a point) like

$$\bar{B}\Gamma_\mu B\bar{B}\tilde{\Gamma}B \tag{4}$$

with $\tilde{\Gamma}$ and Γ_μ representing some Lorentz scalar and Lorentz vector quantities consistent with chiral symmetry, respectively, are subdominant in the chiral counting and can be ignored. Now terms like (4) show up in the chiral Lagrangian as a result of integrating out heavy degrees of freedom with an energy scale $E \gtrsim \Lambda_\chi$. Consider for instance the scalar meson σ that plays an important role in effective field theory of nuclei, a prototype of which being the Walecka model. In free space, there is no low-lying scalar that can appear in the low-energy chiral Lagrangian. But there is a high-lying scalar field that can be associated

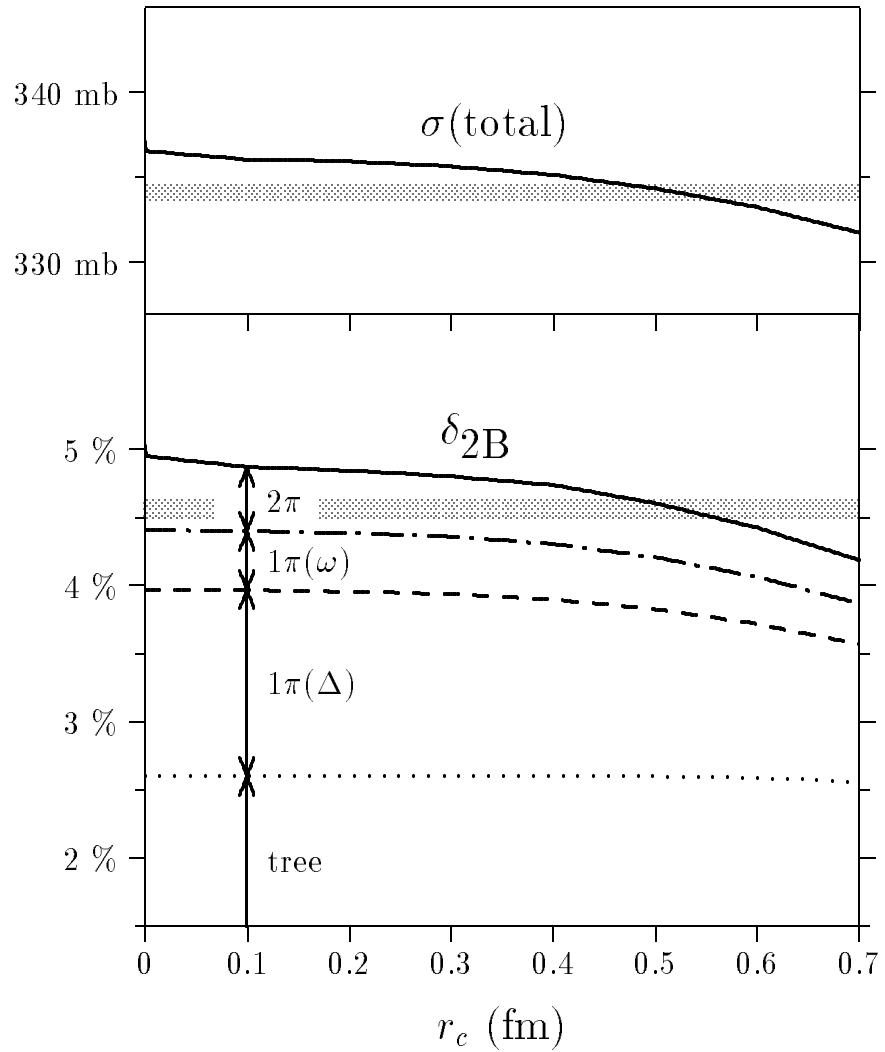


Figure 1: The np capture rate calculated in chiral perturbation theory. The predicted capture rate (upper plot) $\sigma = (334 \pm 3)$ mb agrees with the experimental value (334.2 ± 0.5) mb. The lower plot shows the ratios of the matrix elements of the two-body terms over the one-body term. “Tree” corresponds to the leading chiral order one-pion exchange term (with the blob in Figs.2(a,b) replaced by a bare vertex), “ $1\pi(\omega)$ ” and “ $1\pi(\Delta)$ ” correspond to the next-to-leading order corrections involving the ω meson and Δ resonance, respectively, to the one-pion exchange tree terms Figs.2(a,b). “ 2π ” is the genuine loop correction to the tree contribution. The hadrons appearing in this calculation all have free-space properties.

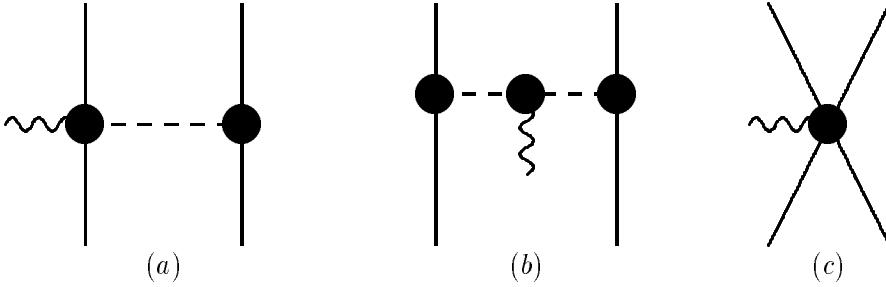


Figure 2: Generic graphs contributing to exchange currents. (a) and (b) are one-pion exchange and (c) represents multipion and/or heavy-meson exchange currents. The large filled circles represent one-nucleon, one-pion irreducible graphs, the solid line the nucleon, the dotted line the pion and the wiggly line the current j_μ .

with the trace anomaly of QCD

$$(T_\mu^\mu)_{QCD} = -\frac{\beta(g)}{2g}(G_{\mu\nu}^a)^2 \sim \chi^4 \quad (5)$$

where β is the QCD beta function, g the color gauge coupling constant, $G_{\mu\nu}^a$ the gluon field tensor and χ the scalar glueball field. Now the χ field is massive, with $m_\chi \sim 2$ GeV, so this degree of freedom appears only in the counter terms. It will give rise to a term like

$$\kappa \bar{B} j_\mu^{(1)} B \bar{B} B \quad (6)$$

with the coefficient κ suppressed by the power $(Q/m_\chi)^2$. In the process (1), Q is of order of 40 MeV, so $(Q/m_\chi)^2 < 10^{-3}$. Stated differently an effective two-body term like (6) will be screened by the short-range correlation implicit in the wave functions.

The situation is quite different, however, in dense nuclear medium. As discussed in [8], as density increases, the scalar (or precisely the quarkish component of the scalar) χ moves downwards in energy and at some high density, it joins the triplet of nearly massless pions to make up the quartet of the $O(4)$ symmetry of chiral symmetry. The merging presumably takes place at the chiral transition point discussed below. The point is that as discussed by Beane and van Kolck [9], in order to reconcile Weinberg's "mended symmetry" [10] with effective chiral Lagrangians at some shorter length scale, the scalar field must come down as a dilaton. Now if it comes down below the chiral scale Λ_χ , then we can no longer consider the scalar as a counter-term contribution. It will strongly couple to low-mass multipion excitations giving among others what is usually taken as a scalar field σ in effective nuclear forces [8]. All other hadrons (other than Goldstone bosons) will also couple to this scalar as well and will undergo a mass shift as density increases. A chiral Lagrangian that accounts for this phenomenon has been shown to lead to the Brown-Rho scaling [11]

$$m_B^*/m_B \approx m_V^*/m_V \approx m_\sigma^*/m_\sigma \approx f_\pi^*/f_\pi \approx \dots \quad (7)$$

where B stands for baryons, V for vector mesons and σ for the dilatonic scalar. The star stands for density-dependent quantities. Such a chiral Lagrangian effective in dense system will then contain these effective constants instead of the free-space values used in the process (1) while preserving the free-space chiral symmetry. As I will argue, the consequence of this scaling can be significant in heavy nuclei and nuclear matter.

One can already see the effect of this scaling in finite nuclei. One clear case is the axial charge transition in nuclei. It was shown in [5, 12] that to order Q^3 in chiral expansion the axial charge transition matrix element in heavy nuclei

$$A(0^\pm) \rightarrow B(0^\mp), \quad \Delta T = 1 \quad (8)$$

is enhanced with respect to the impulse approximation by the factor

$$\epsilon_{MEC} = \frac{m_N}{m_N^*}(1 + R) \quad (9)$$

where R is the ratio of the exchange current matrix element Fig.2 to the impulse approximation calculated with matter-free-space constants. In (8), R is essentially given by Fig.2(a) with the bare coupling and is given by $R \approx 0.5$ with a small variation with density of the system. Now at nuclear matter density, $m_N^*/m_N \approx 0.75$, so

$$\epsilon_{MEC} \approx 2. \quad (10)$$

In light nuclei, we expect $\epsilon_{MEC} \gtrsim 1.5$. These results are in agreement what was found experimentally [13].

One predicts a similar effect in magnetic moments of heavy nuclei but here one has to include other effects of equal importance present in the vector current case. For instance, the “back-flow” correction due to Galilean invariance cancels almost completely the corrections coming from the scaled nucleon mass.

We now turn to an important issue of making a bridge between the chiral theory and Walecka mean field theory of nuclei and nuclear matter which is found to be very successful. I wish to show here, following [14], that Walecka theory is equivalent to the chiral Lagrangian theory at mean field *with* the BR scaling and that this would allow a treatment of fluctuations into different flavor directions (such as strangeness) in a way consistent with the properties of normal nuclear matter. To do this we can focus on the four-fermi interactions allowed in chiral Lagrangians that are relevant in making contact with Walecka theory,

$$\mathcal{L}_{4f} = \alpha (\bar{B}B)^2 + \beta (\bar{B}v^\mu B)^2 \quad (11)$$

where α and β are dimension -2 constants and we are using the heavy-fermion formalism so that v_μ is the velocity four-vector of the heavy baryon. Let us imagine that the first term

of (11) arises from integrating out the heavy chiral singlet scalar χ and the second term from integrating out a heavy chiral singlet vector meson ω . We can include other degrees of freedom in a similar way but we will not need them for symmetric nuclear matter that we shall consider. In this case, we can identify the constants

$$\alpha = \frac{g_\chi^2}{2m_\chi^2}, \quad \beta = -\frac{g_\omega^2}{2m_\omega^2}. \quad (12)$$

We now ask what happens to this Lagrangian when it is immersed in dense and/or hot matter. In mean field, we get the nucleon scalar potential S_N and vector potential V_N as

$$S_N = -\frac{g_\sigma^{*\!2}}{m_\sigma^{*\!2}}\rho_s, \quad (13)$$

$$V_N = \frac{9}{8f_\pi^{*\!2}}\rho \quad (14)$$

where ρ_s is the scalar density and ρ the vector density. In obtaining (14), we have used $SU(3)_f$ relations together with KSRF relation which is known to hold well and put the stars in (13) and (14) to indicate that they are in-medium quantities. Now comparing with the phenomenology with Walecka model, we find that the identification requires that (7) holds with

$$\frac{f_\pi^*}{f_\pi} \approx 0.77. \quad (15)$$

Together with what we found in the case of the axial charge transition, we come to the conclusion that the mean field Walecka theory is just the mean field chiral Lagrangian theory *with BR scaling*.

It is a well-known defect of the mean field Walecka model that the compressibility modulus K_0 is much too high in the model. Now how does this defect get rectified? The answer must lie in higher loop corrections going beyond the mean field as the scaling is known to fail to give the nuclear matter saturation [15]. This is also seen in recent work of Furnstahl et al [16] who note that by giving an anomalous dimension 2.7 to the scalar field σ with the Lagrangian suitably implemented with the trace anomaly, they can obtain the low $K_0 \approx 200$ MeV and the suppression of the many-body terms σ^n , $n > 2$. It is plausible that the anomalous dimension is mocking up the quantum loop effects that seem to be needed in the mean field approach given in [14]. A remarkable observation is that at the anomalous dimension of $d_a \approx 2.7$, two things happen simultaneously. One is that the K_0 which is large at smaller anomalous dimensions stabilizes at ~ 200 MeV for $d_a \approx 2.7$, stays at that value for higher d_a 's and secondly, it is at this fine-tuned value of d_a that *all* multi-body forces get suppressed. This clearly calls for a simple explanation[17].

The above result immediately suggests how to calculate kaon-nuclear interactions in consistency with the nuclear matter properties as given by Walecka theory. To see this,

consider a part of the chiral Lagrangian that figures importantly in the kaon-nuclear sector

$$\mathcal{L}_{KN} = \frac{-6i}{8f^2} (\bar{B}\gamma_0 B) \bar{K} \partial_t K + \frac{\Sigma_{KN}}{f^2} (\bar{B}B) \bar{K} K \equiv \mathcal{L}_\omega + \mathcal{L}_\sigma \quad (16)$$

where $K^T = (K^+ K^0)$. The constant f in (16) can be identified in free space with the pion decay constant f_π . In medium, however, it can be modified as we shall see shortly. In chiral perturbation expansion, the first term corresponds to $\mathcal{O}(Q)$ and the second term to $\mathcal{O}(Q^2)$. There is one more $\mathcal{O}(Q^2)$ term proportional to ∂_t^2 which will be taken into account in the numerical results quoted below but they are not important except for quantitative details.

One can interpret the first term of (16) as arising from integrating out the ω meson as in the baryon sector. The resulting $K^- N$ vector potential in medium can then be deduced in the same way as for V_N :

$$V_{K^\pm} = \pm \frac{3}{8f_\pi^{*2}} \rho. \quad (17)$$

Thus in medium, we may set $f \approx f_\pi^*$ and obtain

$$V_{K^\pm} = \pm \frac{1}{3} V_N. \quad (18)$$

This just says that the ω couples to a *matter field* kaon, hence 1/3 of the ω coupling to the nucleon. The reason for this matter-field nature of the kaon is that all nonstrange hadrons become light in dense medium, so the kaon becomes in some sense heavy. This dual character is known from the hyperon structure which is well described by considering the kaon to be heavy as in the Callan-Klebanov model.

As for the second term of (16), we use that the kaon behaves as a massive matter field. We therefore expect that it be coupled to the chiral scalar χ as

$$\mathcal{L}_\sigma = \frac{1}{3} 2m_K g_\sigma^* \bar{K} K \chi \quad (19)$$

where the factor 1/3 accounts for one non-strange quark in the kaon as compared with three in the nucleon. When the χ field is integrated out as above, we will get, analogously to the nucleon case,

$$\mathcal{L}_\sigma = 2m_K \frac{1}{3} \frac{g_\sigma^{*2}}{m_\sigma^{*2}} \bar{B} B \bar{K} K. \quad (20)$$

Comparing with the second term of (16), we find

$$\frac{\Sigma_{KN}}{f^2} \approx 2 \frac{m_K}{3} \frac{g_\sigma^{*2}}{m_\sigma^{*2}}. \quad (21)$$

We can get the Σ_{KN} from lattice calculations [18], $\Sigma_{KN} \approx 3.2m_\pi$. This gives $f \approx f_\pi^*$. Therefore we have

$$S_{K^\pm} = \frac{1}{3} S_N. \quad (22)$$

To summarize: the kaon-nuclear potential gotten from a chiral Lagrangian and the nucleon-nuclear potential given by Walecka mean field theory are directly related through BR scaling.

Given Walecka mean fields for nucleons, we can now calculate the corresponding mean-field potential for K^- -nuclear interactions in symmetric nuclear matter. From the results obtained above, we have

$$S_{K^-} + V_{K^-} \approx \frac{1}{3}(S_N - V_N). \quad (23)$$

Phenomenology in Walecka mean-field theory gives $(S_N - V_N) \lesssim -600$ MeV for $\rho = \rho_0$. This leads to the prediction that at nuclear matter density

$$S_{K^-} + V_{K^-} \lesssim -200 \text{ MeV}. \quad (24)$$

This seems to be consistent with the result of the analysis in K-mesic atoms made by Friedman, Gal and Batty [19] who find attraction at $\rho \approx 0.97\rho_0$ of

$$S_{K^-} + V_{K^-} = -200 \pm 20 \text{ MeV}. \quad (25)$$

An immediate consequence of this mean field description of the kaonic sector is that kaons will condense in dense neutron star (or nuclear star) matter at a density

$$\rho_c \sim 2\rho \quad (26)$$

as found by Lee et al [20] in ChPT to one-loop order.

To conclude, I make a few remarks on the nature of the scaling properties.

- There are two points to the issue. One is that *some or all* light hadrons may be undergoing a downward mass shift as density or temperature is increased. The second is that the scaling is like (7). These are basically two different issues. To the extent that hadron masses are generated spontaneously, it is inevitable that at least some masses should drop. Indeed QCD sum rule calculations do predict the drop for the vector-meson mass [21], the most recent value being [22]

$$m_\rho^*/m_\rho = 0.78 \pm 0.08, \quad \rho \approx \rho_0. \quad (27)$$

On the other hand, the “universal scaling” (7), obtained at mean field, may not be strictly valid. In fact, large N_c arguments imply $m_N^*/m_N \approx \sqrt{g_A^*/g_A}(f_\pi^*/f_\pi)$. If we look at the strict $N_c = \infty$ limit, $g_A^* = g_A$ as we showed in [11], the scaling (7) holds but we know that in nuclei, $g_A^* \approx 1$ and thus in finite nuclei we expect that the nucleon scales somewhat faster than the pion decay constant, at least up to $\rho \sim \rho_0$. Whether the “universal scaling” (7) is consistent with nature remains to be seen.

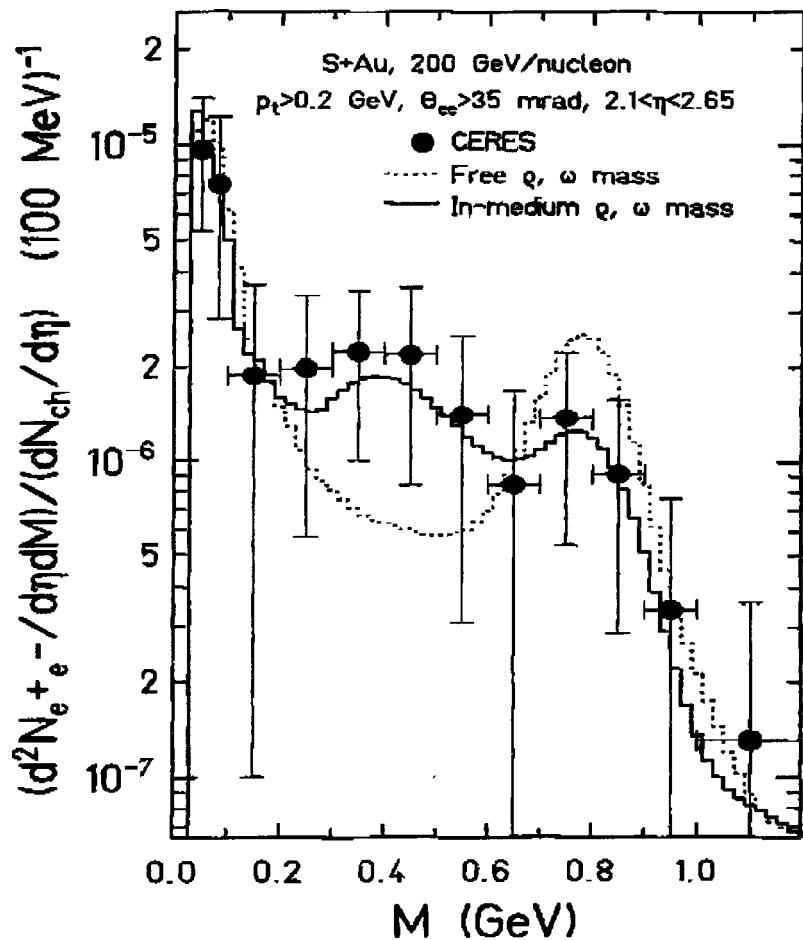


Figure 3: The Li-Ko-Brown explanation of the dilepton data of the CERES collaboration. The dotted line is the theoretical prediction *without* the scaling of the ρ and ω mesons and the solid curve *with* the BR scaling. The three lowest mass points are essentially given by Dalitz pairs, so the relevant data points are the ones for higher invariant masses. Note that the peak shown at $M \sim 800$ MeV is predominantly given by the ω decay outside of the medium.

2. An intriguing question is how far the scaling (7) can be pushed in density and/or temperature. Can one use the mean field argument all the way to the chiral phase transition?

This is a highly relevant question since there are arguments [23] that the second order chiral phase transition relevant to QCD with 2 flavors is of mean field type as in 3D Gross-Neveu model. This suggests that the mean field chiral theory or Walecka theory could be used to discuss the QCD chiral phase transition [24]. If correct, this theory will be very useful for studying phase transitions in heavy ion collisions.

3. There are a large number of experimental projects to measure the mass shift of hadrons in dense matter, particularly at GSI, CEBAF and CERN using dilepton-pair production. Recent data from the CERES collaboration [25] on e^+e^- pairs in S on Au collisions at 200 GeV/nucleon can be understood in terms of the scaled mass of the vector mesons in medium, primarily due to density effect. In a recent paper, Li, Ko and Brown [26] have shown that the enhancement of the produced lepton pairs observed in the range of invariant mass $300 \text{ MeV} \lesssim M \lesssim 550 \text{ MeV}$ can be explained simply by the BR scaling in the ρ mass as the pairs are produced mainly through $\pi^+\pi^- \rightarrow \rho^* \rightarrow e^+e^-$. The fit is given in Fig.3. The analysis is a complex one involving the assumption of an expanding fireball in chemical equilibrium, but the economy of the explanation and the quantitative success make it quite compelling.

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